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QUANTUM 1/f NOISE IN HIGH TECHNOLOGY APPLICATIONS INCLUDING ULTRASMALL STRUCTURES AND DEVICES

THIRD ANNUAL REPORT

June 15, 1991 - June 14, 1992

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Abstract

The practical application of the quantum 1/f effect to quartz resonators and to infrared detectors present in this report allows us for the first time to understand and to extend the stability limits of quartz resonators. It also explains 1/f noise in most semiconductor devices and in infrared detectors considered in this report in the presence of radiation, although 1/f noise has been already successfully pushed below background in PtSi on p-type Si Schottky diodes at RADC/Hanscomb. A fundamental breakthrough was performed through the first direct derivation of the coherent quantum 1/f effect from a special quantum-electrodynamic propagator, and from the author's general sufficient 1/f chaos criterion presented in the previous yearly report. Finally, the quantum 1/f cross-correlations derived by the author have been used to recalculate and to graph the quantum 1/f mobility fluctuations in Si and GaAs samples as a function of temperature and doping, in good agreement with the measurements of Tacano in Japan and Hooge in the Netherlands.

CONTENTS

I. INTRODUCTION

The progress achieved this year includes the derivation of the Coherent Quantum 1/f Effect from a special quantum-electrodynamical propagator known since 1975, and from the sufficient 1/f criterion introduced in the previous Annual Report (Sec. II and III, respectively). It also includes the successful practical application of the quantum 1/f theory to quartz resonators and the verification of the results at NIST-Boulder with the cooperation of Fred Walls (Sec. IV). It further includes in Sec. V the practical application of the quantum 1/f theory to infrared detectors in the presence of radiation, in particular, a proof of the absence of quantum 1/f noise in the process of carrier photogeneration in semiconductors. it finally includes in Sec. VI the recalculation and first time graphical representation of quantum 1/f mobility fluctuations in Si and GaAs samples, based on the author's new cross-correlation formula, as a function of temperature and doping level.

We also mention the NSF-sponsored organization by the author of the "V. International van der Ziel Symposium on Quantum 1/f Noise and Other Low-Frequency Fluctuations" at the University of Missouri-St. Louis on May 22-23, 1992 with a 100% larger participation than ever before, and the ongoing organization by the author of the "XII Int. Conf. on Noise in Physical Systems and 1/f Noise" in St. Lois, for Aug. 16-20, 1993. Finally, we mention the creation of an Institute of Molecular Electronics at the University of Missouri St. Louis with the author's participation in September 1991. My collaborators have been F. Walls, E. Bernardi, T. Chung, A. Först-Chung, Y. Zhang, Xuewei Hu, Jian Xu, and I. Proleiko. These results will be briefly presented below.

II. COHERENT QUANTUM 1/f CHAOS

Conventional quantum 1/f fluctuations of physical cross sections and process rates have been introduced by us as a fundamental infrared divergence phenomenon in 1975 [1]. Some of the subsequent publications [2]-[14] have shown this new effect to be unaffected by the presence of the thermal radiation background [4], [5], some have derived it with wave packets [11], including a finite mean free path [9], in second quantization [13], with the Keldysh-Schwinger method [12], or in the Van Hove weak interaction limit [14], some derived its characteristic functional [7], or applied it to the calculation [10] of mobility and recombination speed fluctuations in semiconductors and semiconductor devices. Others [15] - [27] verified the new effect experimentally and successfully applied it to electronic devices.

The present paper derives a related fundamental effect which we call the coherent 1/f effect, with elementary methods of quantum electrodynamics and non-relativistic many-body theory. Our derivation uses the new picture introduced by Dollard [28] and generalized by

Kulish and Fadde'ev [29] and later by Zwanziger [30], in agreement with earlier work by Chung [31] and by Kibble [32]. In this new picture, the asymptotic Coulomb interaction is included in the unperturbed Hamiltonian rather than in the perturbation part. This leads to a more complex physical free particle notion which includes a coherent photon cloud, and replaces the pole in the propagator with a branch point. It also leads to a smeared-out mass shell. Using this picture, we can neglect the remaining part of the interaction if we limit ourselves to the asymptotic region of large distances and times, which are important in the case of 1/f noise.

For N electrons in a Fermi sphere shifted in momentum space by a vector p_o and occupying N/2 orbitals e^{ipr}, the propagator derived by these authors [30] can be reduced for large time components of x'-x to the non-relativistic form

$$\begin{aligned} -i\langle\Phi_o|T\psi_{s'}(x')\psi_s^\dagger(x)|\Phi_o\rangle &= \delta_{ss'}\, G_s(x'-x) \\ &= (i/V)\sum_p \{\exp i[p(r-r')-p^2(t-t')/2m]/\hbar\} n_{p,s} \\ &\quad \times\{-ip(r-r')/\hbar+i(m^2c^2+p^2)^{1/2}(t-t')(c/\hbar)\}^{\alpha/\pi}. \end{aligned} \qquad (1)$$

Here $\alpha=e^2/\hbar c=1/137$ is Sommerfeld's fine structure constant, $n_{p,s}$ the number of electrons in the state of momentum p and spin s, m the rest mass of the fermions, $\delta_{ss'}$ the Kronecker symbol, c the speed of light, x=(r,t) any space-time point and V the volume of a normalization box. T is the time-ordering operator which orders the operators in the order of decreasing times from left to right and multiplies the result by $(-1)^P$, where P is the parity of the permutation required to achieve this order. For equal times, T normal-orders the operators, i.e., for t=t' the left-hand side of Eq. (1) is $i\langle\Phi_o|\psi_s^\dagger(x)\psi_{s'}(x')|\Phi_o\rangle$. The state Φ_o of the N electrons is described by a Slater determinant of single-particle orbitals.

Consider first the case t=t' for simplicity, although only the case of large t-t' can be expected to be experimentally applicable. The pair correlation function can then be decomposed as follows

$$\begin{aligned} \langle\Phi_o|\psi_s^\dagger(x)\psi_{s'}^\dagger(x')\psi_{s'}(x')\psi_s(x)|\Phi_o\rangle &= \langle\Phi_o|\psi_s^\dagger(x)\psi_s(x)|\Phi_o\rangle\langle\Phi_o|\psi_{s'}^\dagger(x')\psi_{s'}(x')|\Phi_o\rangle \\ &\quad - \langle\Phi_o|\psi_s^\dagger(x)\psi_{s'}(x')|\Phi_o\rangle\langle\Phi_o|\psi_{s'}^\dagger(x')\psi_s(x)|\Phi_o\rangle. \end{aligned} \qquad (2)$$

The first term can be expressed in terms of the particle density of spin s, n/2 = N/2V = $\langle\Phi_o|\psi_s^\dagger(x)\psi_s(x)|\Phi_o\rangle$, while the second term can be expressed in terms of the Green function (1) in the form

$$A_{ss'}(x-x') = \langle\Phi_o|\psi_s^\dagger(x)\psi_{s'}^\dagger(x')\psi_{s'}(x')\psi_s(x)|\Phi_o\rangle = (n/2)^2 + \delta_{ss'}\, G_s(x'-x)G_s(x-x'). \qquad (3)$$

The "relative" autocorrelation function $A(x-x')$ describing the normalized pair correlation independent of spin is obtained by dividing by n^2 and summing over s and s'

$$A(x-x') = 1 + (1/n^2)\sum_{s} G_s(x-x')G_s(x'-x)$$

$$= 1 - (1/N^2)\sum_{s} \sum_{pp'}\{\exp[i(p-p')(r-r')/\hbar]\}n_{p,s}n_{p',s} \times |p(r-r')/\hbar|^{\alpha/\pi}|p'(r-r')/\hbar|^{\alpha/\pi}. \qquad (4)$$

Here we have used Eq. (1). The low-wavenumber part A_l of this relative density autocorrelation function is given by the terms with $p=p'$.

$$A_l(x-x') = 1 - (1/N^2)\sum_{s}\sum_{p} n_{p,s}|p(r-r')/\hbar|^{2\alpha/\pi} \qquad (5)$$

$$= 1 - (2/N^2)[V/(2\pi\hbar)^3] \int_{p<p_F} d^3p|(p+p_o)(r-r')/\hbar|^{2\alpha/\pi}$$

$$= 1 - (2/N^2)[V/(2\pi\hbar)^3] \int_{-a}^{a}dp_1 \int_{-a}^{a}dp_2 \int_{-a}^{a}dp_3|(p_3+p_{o3})(r-r')/\hbar|^{2\alpha/\pi}$$

$$= 1 - (2/N^2)[V/(2\pi\hbar)^3]8a^3|a(r-r')/\hbar|^{2\alpha/\pi}$$

$$= 1 - (1/N)[(\pi/6)^{1/3}p_F|r-r'|/\hbar]^{2\alpha/\pi}. \qquad \text{for } p_F>>p_{o3}; \qquad (6)$$

and

$$= 1 - (1/N)|p_o(r-r')/\hbar|^{2\alpha/\pi} \qquad \text{for } p_F<<p_{o3}; \qquad (6')$$

In view of the smallness of $2\alpha/\pi$, for $p_F>>p_{o3}$ we have integrated in cartesian coordinates, approximating the Fermi sphere by a cube of side 2a with $a= (\pi/6)^{1/3}p_F$. The result is practically independent of p_o and of $p_{o3}=p_o(r-r')/|r-r'|$. The factor $(\pi/6)^{2\alpha/3\pi}$ can be neglected. For $p_F<<p_{o3}$ we used the mean value theorem for estimating the integral over d^3p in spherical coordinates. In all cases the autocorrelation decreases very slowly from 1 when $|r-r'|$ is increased to very large values. Writing the rectangular bracket in Eq. (6) as an exponential function of its logarithm, expanding the resulting exponential, and keeping only the first term, we obtain with $p_F/\hbar=k_F$

$$A_l(x-x') = (1-2/N) + (6/\pi)^{2\alpha/3\pi}\{[1/k_F|r-r'|]^{2\alpha/\pi}\}/N$$

$$= (1-2/N) + (6/\pi)^{2\alpha/3\pi}[1/k_F]^{2\alpha/\pi}(2\alpha/N\pi)\int_0^\infty \cos[k|r-r'|]dk/k^{1-2\alpha/\pi}.$$

$$\approx\{N-2 + (2\alpha/\pi)\int_0^\infty [k/k_F]^{2\alpha/\pi}\cos[k|r-r'|]dk/k\}/N \quad \text{for } p_F>>p_{o3}; \qquad (7)$$

$$\approx\{N-2 + (2\alpha/\pi)\int_0^\infty [k/k_o]^{2\alpha/\pi}\cos[kp_o(r-r')/p_o]\, dk/k\}/N \quad \text{for } p_F<<p_{o3}; \qquad (7')$$

Here we have used a well-known Fourier integral [35] and we have introduced $k_o=p_{o3}/\hbar$. According to the Wiener-Khintchine theorem, the coefficient of the cos gives the spectral density. To get it for the fractional fluctuations $\delta n/n$, we divide by the constant term N-2

$$S_{\delta n/n}(k) = [2\alpha/\pi k(N-2)][k/K]^{2\alpha/\pi}, \qquad (8)$$

where $K=k_F$ for $p_F>>p_{o3}$ as in the case of metals with spherical wave symmetry, and $K=k_o$ for $p_F<<p_{o3}$. Although inapplicable, this pure 1/k spectrum is the wave-number equivalent of the coherent quantum 1/f noise derived earlier [33], [34] in excellent agreement with the experiments on large electronic devices [15], [25]. Due to $2\alpha A<<1$ the second factor is practically unity and of no importance, except for eliminating the logarithmic divergence from the spectral integral. This wave number spectrum also entails a 1/f frequency spectrum obtained by writing dk/k=df/f, as was shown in detail in a previous paper [13]. For equal times our result can not be expected to be valid, due to the asymptotic character of Eq. (1). We shall now derive the 1/f spectrum directly below.

If $t\neq t'$, Eq. (2) is replaced by

$$\langle\Phi_o|T\psi_s^\dagger(x)\psi_s(x)\psi_{s'}^\dagger(x')\psi_{s'}(x')|\Phi_o\rangle = \langle\Phi_o|\psi_s^\dagger(x)\psi_s(x)|\Phi_o\rangle\langle\Phi_o|\psi_{s'}^\dagger(x')\psi_{s'}(x')|\Phi_o\rangle$$
$$- \langle\Phi_o|T\psi_{s'}(x')\psi_s^\dagger(x)|\Phi_o\rangle\langle\Phi_o|T\psi_s(x)\psi_{s'}^\dagger(x')|\Phi_o\rangle. \qquad (9)$$

Eq. (3) remains the same, except for the middle part which is replaced by the left hand side of Eq. (9). Eq. (4) becomes now

$$A(x-x') = 1 - (1/n^2)\sum_s G_s(x-x')G_s(x'-x)$$

$$= 1 - (1/N^2)\sum_s \sum_{pp'}\{\exp i[(p-p')(r-r')-(p^2-p'^2)(t-t')/2m]/\hbar\}n_{p,s}n_{p',s}$$

$$\times\{p(r-r')/\hbar-(m^2c^2+p^2)^{1/2}(t-t')(c/\hbar)\}^{\alpha/\pi}$$
$$\times\{p'(r-r')/\hbar-(m^2c^2+p'^2)^{1/2}(t-t')(c/\hbar)\}^{\alpha/\pi}. \tag{10}$$

Here we have used again Eq. (1). The low-frequency and low-wavenumber part A_l of this relative density autocorrelation function is also given by the terms with $p=p'$.

$$A_l(x-x') = 1 - (1/N^2)\sum_s \sum_p n_{p,s}$$
$$\times\{p(r-r')/\hbar-(m^2c^2+p^2)^{1/2}(t-t')(c/\hbar)\}^{2\alpha/\pi} \tag{11}$$

$$=1 -(2/N^2)[V/(2\pi\hbar)^3 \int_{p<p_F} d^3p|(p+p_o)(r-r')/\hbar-[m^2c^2+(p+p_o)^2]^{1/2}(t-t')(c/\hbar)|^{2\alpha/\pi}$$

$$= 1 - (1/N)|p_o(r-r')/\hbar-mc^2\tau/\hbar|^{2\alpha/\pi} \qquad \text{for } p_F<<|p_{o3}-mc^2\tau/z|. \tag{12}$$

Here we have used the mean value theorem, considering the $2\alpha/\pi$ power as a slowly varying function of p and neglecting p_o in the coefficient of $\tau=t-t'$, with $z=|r-r'|$. Writing the power again as an exponential function of its logarithm, expanding the resulting exponential, and keeping only the first term, we obtain with $(\hbar/mc^2)|p_o(r-r')/\hbar-mc^2\tau/\hbar|=\theta=|\tau-p_o(r-r')/mc^2|$

$$A_l(x-x') = 1-[(mc^2/\hbar)\theta]^{2\alpha/\pi}/N$$
$$= (1-2/N) + [\hbar/mc^2]^{2\alpha/\pi}(2\alpha/N\pi)\int_0^\infty \cos[\omega\theta]d\omega/\omega^{1-2\alpha/\pi}.$$
$$= \{N-2 + (2\alpha/\pi)\int_0^\infty [\hbar\omega/mc^2]^{2\alpha/\pi}\cos[\omega\theta]\, d\omega/\omega\}/N. \tag{13}$$

According to the Wiener-Khintchine theorem, the coefficient of the cos gives the spectral density. To get it for the fractional fluctuations $\delta n/n$, we divide by the constant term N-2

$$S_{\delta n/n}(k) = [2\alpha/\pi\omega(N-2)][\hbar\omega/mc^2]^{2\alpha/\pi}. \tag{14}$$

This result is also applicable for the particle current fluctuation spectrum. Indeed, for current density fluctuations δj we include a $(\hbar/mi)\nabla$ in front of each of the two ψ operators in Eq. (9), a factor pp'/p_o^2 in Eq. (10) after the summation signs, a factor $(p/p_o)^2$ in the first form of Eq.

(11), a factor $(p+p_o)^2/p_o^2$ in the second form, and no changes in Eqs. (12)-(13). Eq. (14) becomes

$$\underline{S_{\delta j/j}(k) = [2\alpha/\pi\omega(N-2)][\hbar\omega/mc^2]^{2\alpha/\pi}}. \tag{15}$$

This result coincides with our earlier theoretical result for coherent quantum 1/f noise if we replace N with N-2. The validity of this equation is restricted to low frequencies and wave-numbers. This equation is in excellent agreement with mobility and diffusion 1/f noise in large devices.

Finally, we consider the errors caused by the neglect of higher order terms in the expansion of the exponential functions resulting from Eqs. (6)-(6') and (12). For a thermal electron and r=1cm in Eq. (6') we get $kr=10^{12}$ and $(2\alpha/\pi)\ln(kr)=0.12$, yielding an error of 12%. For $t=10^{17}$s in Eq. (12), which is the age of the universe, we get an error of 40% as the upper limit. Therefore a more exact treatment is of some interest. Using the identity [35]

$$\theta^{2\alpha/\pi} = [-(2\alpha/\pi)\int_{\omega_o}^{\infty}\omega^{-2\alpha/\pi}\cos(\theta\omega)d\omega/\omega]$$
$$\times\{\cos\alpha + (2\alpha/\pi)\sum_{n=0}^{\infty}(\theta\omega_o)^{2n-2\alpha/\pi}[(2n)!(2n-2\alpha/\pi)]^{-1}\}^{-1}, \tag{16}$$

with arbitrarily small cutoff ω_o, we obtain from Eq. (12) the exact form

$$A_j(x-x') = 1 + [(2\alpha/\pi N)\int_{\omega_o}^{\infty}(mc^2/\hbar\omega)^{2\alpha/\pi}\cos(\theta\omega)d\omega/\omega]$$
$$\times\{\cos\alpha + (2\alpha/\pi)\sum_{n=0}^{\infty}(\theta\omega_o)^{2n-2\alpha/\pi}[(2n)!(2n-2\alpha/\pi)]^{-1}\}^{-1}. \tag{17}$$

This would indicate a $\omega^{-1-2\alpha/\pi}$ spectrum and a 1/N dependence of the spectrum of fractional n and j fluctuations, if we neglect the curly bracket in the denominator which is close to unity for very small ω_o. We thus realize that the unusual N-2 dependence in Eqs. (14)-(15) is caused by the forced introduction of the integrable $\omega^{2\alpha/\pi-1}$ spectrum in place of the $\omega^{-1-2\alpha/\pi}$ spectrum. Due to the smallness of α both forms coincide in practical applications. Eq. (15) for the coherent QED chaos process in electric currents can thus be written also in the form

$$\underline{S_{\delta j/j}(k) = [2\alpha/\pi\omega N][mc^2/\hbar\omega]^{2\alpha/\pi} = 2\alpha/\pi\omega N = 0.00465/\omega N}. \tag{18}$$

This result derived directly earlier [33], [34], is in excellent agreement with the measurements [15], [25], in large {see [34] for a definition of "large" or "extended", and for an interpolation with the conventional quantum 1/f chaos effect} devices such as large n^+p $Hg_{1-x}Cd_xTe$ infrared detector diodes. It is also close to the empirical value of 0.002/ωN observed earlier by Hooge [36] in semiconductors and metals. Being observed in the presence of a constant applied field, these fundamental quantum current fluctuations are usually interpreted as mobility fluctuations.

III. APPLICATION OF THE SUFFICIENT 1/f CRITERION TO QUANTUM 1/f CHAOS

The nonlinearity causing the 1/f spectrum of turbulence in both semiconductors and metals is caused by the reaction of the field generated by charged particles and their currents back on themselves. The same nonlinearity is present in quantum electrodynamics (QED), where it causes the infrared divergence, the infrared radiative corrections for cross sections and process rates, and the quantum 1/f effect. We shall prove this on the basis of the sufficient criterion for 1/f spectral density in chaotic systems, derived in the previous annual report.

Consider a beam of charged particles propagating in a well-defined direction which we shall call the x direction, so that the one-dimensional Schrödinger equation describes the longitudinal fluctuations in the concentration of particles. Considering the non-relativistic case which is encountered in most quantum 1/f noise applications, we write in second quantization the equation of motion for the Heisenberg field operators ψ of the in the form

$$i\hbar\partial\psi/\partial t = (1/2m)[-i\hbar\nabla - (e/c)\, A]^2\psi, \tag{19}$$

With the non-relativistic form $J = -i\hbar\psi^*\nabla\psi/m$ + hermitic conjugate, and with

$$A(x,y,z,t) = (\hbar/2cmi)\int\frac{[\psi^*\nabla\psi-\psi\nabla\psi^*]}{|x-x'|}dx' \tag{20}$$

we obtain

$$i\hbar\partial\psi/\partial t = (1/2m)\Big[-i\hbar\nabla - (e\hbar/2c^2mi)\int\frac{[\psi^*\nabla\psi-\psi\nabla\psi^*]}{|x-x'|}dx'\Big]^2\psi. \tag{21}$$

At very low frequencies or wavenumbers the last term in rectangular brackets is dominant on the r.h.s., leading to

$$i\hbar\partial\psi/\partial t = (-1/2m)\Big[(e\hbar/2c^2m)\int\frac{[\psi^*\nabla\psi-\psi\nabla\psi^*]}{|x-x'|}dx'\Big]^2\psi. \qquad (22)$$

For x replaced by λx, and x' replaced by $\lambda x'$, we obtain

$$i\hbar\partial\psi/\partial t = (-1/2m)\Big[(e\hbar/2c^2m)\int\frac{[\psi^*\nabla/\lambda\psi-\psi\nabla/\lambda\psi^*]}{\lambda|x-x'|}\lambda^3dx'\Big]^2\psi = \lambda^2H\psi = \lambda^{-p}H\psi. \quad (23)$$

This satisfies our homogeneity criterion with p=-2. Our sufficient criterion only requires homogeneity, with any value of the weight p, for the existence of a 1/f spectrum in chaos. Therefore, we expect a 1/f spectrum of quantum current-fluctuations, i.e., of cross sections and process rates in physics, as derived in detail in Sec. II above. This is in agreement with the well-known, and experimentally verified, results of the Quantum 1/f Theory.

In conclusion, we realize that, both in classical and quantum mechanical nonlinear systems, the limiting behavior at low wave numbers is usually expressed by homogeneous functional dependences, leading to fundamental 1/f spectra on the basis of our criterion.

IV. QUANTUM 1/f FLICKER OF FREQUENCY IN QUARTZ RESONATORS:THEORY AND EXPERIMENT

IV. 1. Introduction

Flicker of frequency noise is an important characteristic of quartz resonators which limits their stability and determines their utility for most applications. The 1/f contribution to frequency stability is best obtained by observing the stability over a range of measurement times in the time domain, or over a range of frequencies in the frequency domain. From the extended data one can fit a flicker of frequency model to the data that excludes the contributions from random walk frequency modulation and from the drift present in many resonators and oscillators (Fig. 1).

One of the first who systematically studied the 1/f noise as a function of geometry, temperature and Q-factors was Gagnepain [37]. He noticed empirically that the 1/f part of the spectral density of fractional frequency fluctuations, $S_y(f)$, varied as Q^{-4} for resonators between 1 and 25 MHz. As the temperature changes, the Q-factor of a resonator changes. This allows us to exclude the effect of many other factors. Additional work by Parker, however, showed that the data from both bulk acoustic wave (BAW) and surface acoustic wave (SAW) devices could be roughly fit to the the same model [38], if one assumes a Q^{-4} dependence for the phase noise S_ϕ insted of S_y. The fit is shown on Fig. 2.

From a theoretical point of view, fundamental work by this author [39], inspired by the quantum 1/f theory [40], has derived the Q^{-4} depenence of S_y for all resonators from first principles already 1979. Work on many systems other than quartz has yielded very good quantitative agreement between theory and experimental data for quantum 1/f noise [41]. This paper refines the previous theoretical work on 1/f frequency noise in quartz to suggest a better framework for predicting the level of 1/f frequency noise in quartz resonators over a wide range of frequencies and Q-factors.

IV. 2. Theory of 1/f Frequency Noise in Quartz Resonators

According to the general quantum 1/f formula [40], $\Gamma^{-2}S_\Gamma(f)=2\alpha A/f$ with $\alpha=e^2/\hbar c=1/137$ and $A=2(\Delta J/ec)^2/3\pi$ is the quantum 1/f effect in any physical process rate Γ. Setting $J=dP/dt=\dot{P}$ where P is the vector of the dipole moment of the quartz crystal, we obtain for the fluctuations in the rate Γ of phonon removal from the main resonator oscillation mode (by scattering n a phonon from any other mode of average frequency $<\omega>$) of the crystal, (or via a two-phonon-process at a crystall defect or impurity, involving a phonon of average frequency $<\omega'>$) the spectral density

$$S_\Gamma(f) = \Gamma^2 4\alpha(\Delta\dot{P})^2/3\pi e^2c^2, \tag{24}$$

where $(\Delta\dot{P})^2$ is the square of the dipole moment rate change associated with the process causing the removal of a phonon from the main oscillator mode. To calculate it, we write the energy W of the interacting resonator mode $<\omega>$ in the form

$$W = n\hbar<\omega> = 2(Nm/2)(dx/dt)^2=(Nm/e^2)(e\ dx/dt)^2=(m/Ne^2)\varepsilon^2(\dot{P})^2; \tag{25}$$

The factor two includes the potential energy contribution. Here m is the reduced mass of the elementary oscillating dipoles, e their charge, ε a polarization costant, and N their number in the quartz crystal. Applying a variation $\Delta n=1$ we get

$$\Delta n/n = 2|\Delta\dot{P}|/|\dot{P}|, \text{ or } \Delta\dot{P}=\dot{P}/2n. \tag{26}$$

Solving Eq. (25) for $\dot{P}$ and substituting, we obtain

$$|\Delta\dot{P}| = (N\hbar<\omega>/n)^{1/2}(e/2\varepsilon) \tag{27}$$

Substituting $\Delta\dot{P}$ into Eq. (24), we get

$$\Gamma^{-2}S_\Gamma(f) = N\alpha\hbar<\omega>/3n\pi mc^2f\varepsilon^2 = A/f . \tag{28}$$

This result is applicable to the fluctuations in the loss rate Γ of the quartz.
The corresponding resonance frequency fluctuations of the quartz resonator are given by[2]

$$\omega^{-2}S_{\omega}(f) = (1/4Q)(\Lambda/f) = N\alpha\hbar<\omega>/12n\pi mc^2f\varepsilon^2Q^4, \quad (29)$$

where Q is the quality factor of the single-mode quartz resonator considered, and $<\omega>$ is not the circular frequency of the main resonator mode, ω_0, but rather the practically constant frequency of the average interacting phonon, considering both three-phonon and two-phonon processes. The corresponding $\Delta\dot{P}$ in the main resonator mode has to be also included in principle, but is negligible because of the very large number of phonons present in the main resonator mode.

Eq. (29) can be written in the form

$$S(f) = \beta V/fQ^4, \quad (30)$$

where, with an intermediary value $<\omega>=10^8$/s, with $n=kT/\hbar<\omega>$, T=300K and $kT=4\ 10^{14}$,

$$\beta = (N/V)\alpha\hbar<\omega>/12n\pi\varepsilon^2mc^2 = 10^{22}(1/137)(10^{-27}10^8)^2/12kT\pi10^{-27}\ 9\ 10^{20} = 1. \quad (31)$$

The form of Eq. (30) shows that the level of 1/f frequency noise depends not only as Q^{-4} as previously proposed, but also on the oscillation frequency or the volume of the active region. This model qualitatively fits the data of Gagnepain [37,48] where he varied the Q-factor with temperature in the same resonator (but not frequency or volume).

The model also provides the basis for predicting how to improve the 1/f level of resonators, beyond just improving the Q-factor, which has been known for many years. Since the level depends on active volume, one should use the lowest overtone and smallest diameter consistent with other circuit parameters.

IV. 3. Experimental Measurements and Analysis of 1/f Noise in Quartz Resonators

The level of 1/f frequency noise in quartz resonators has been measured using phase bridges and complete oscillators [37,38,42-47]. Unfortunately much of the data in the literature is unusable for modeling because the unloaded Q-factor is unknown. (Our case is even more restrictive because we also need to know the electrode size). The phase bridge approach has the advantage that the unloaded Q-factor can be easily measured and the noise in the measurement electronics can be evaluated independent of the resonator. If resonator pairs are used, the noise of the driving source can generally

Phase Noise Model for 5 MHz Oscillator

$$\mathscr{L}(f) = 10^{-12.86}/f^3 + 10^{-15.0}/f + 10^{-178.7}$$

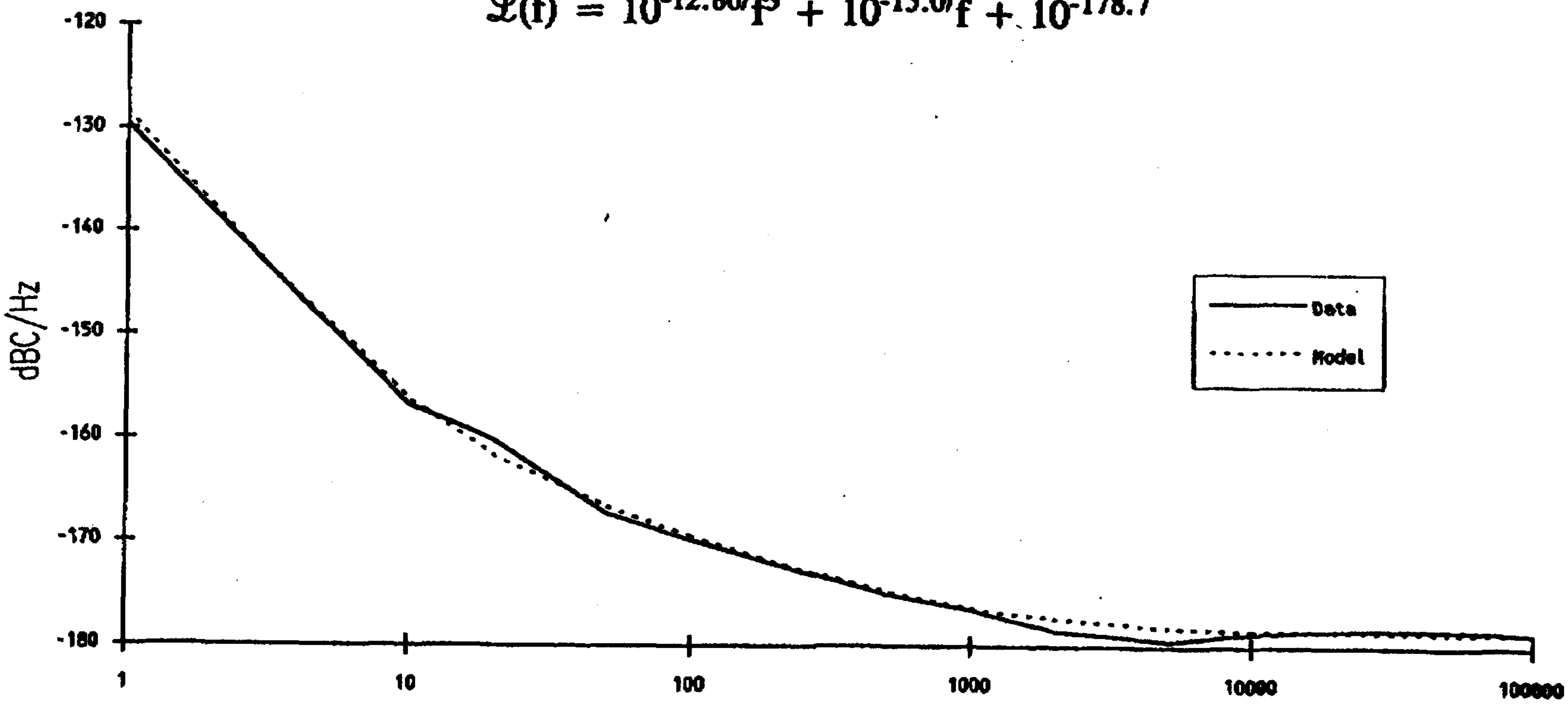

Fig 1

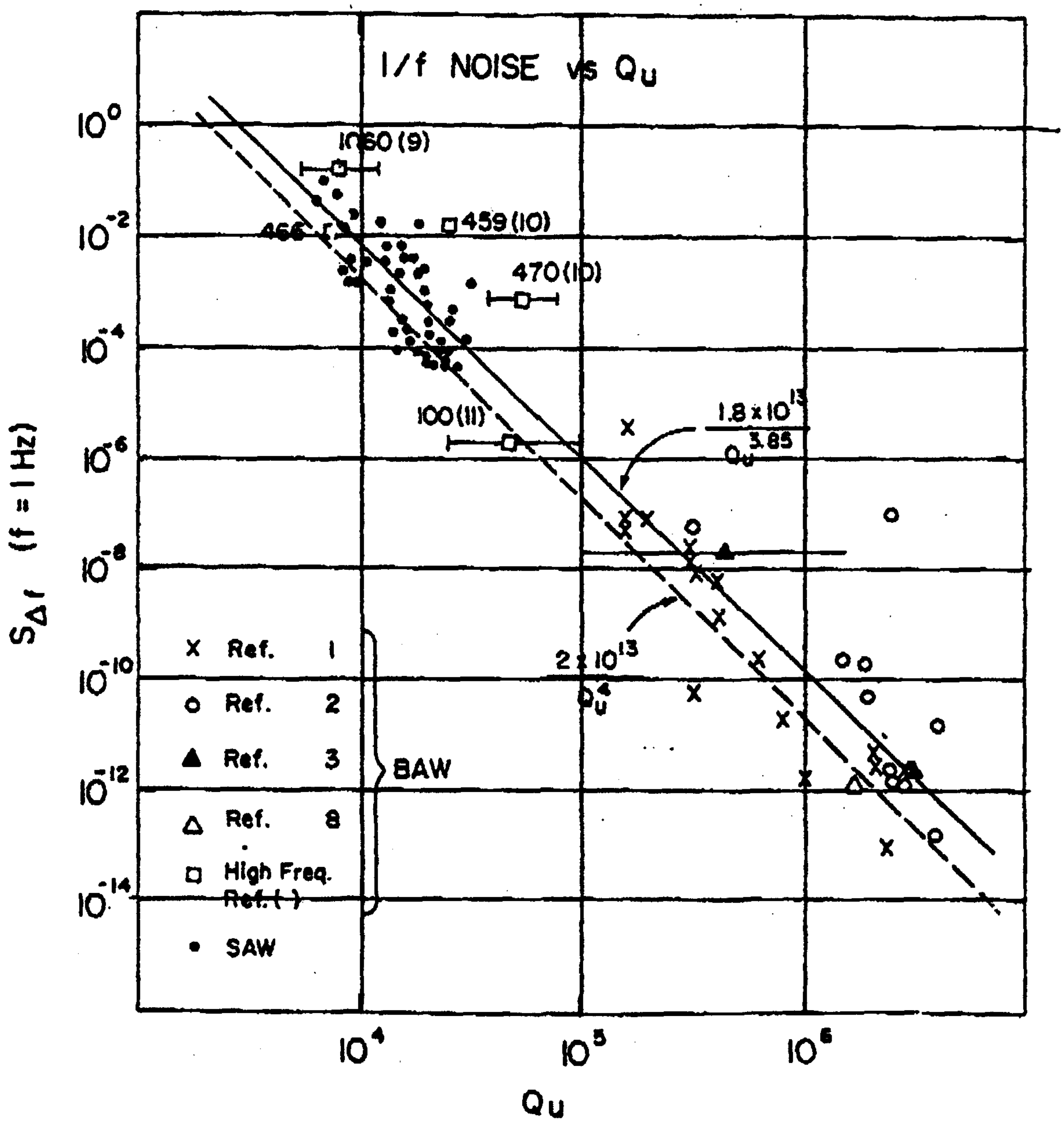

Figure 2 1/f Noise Level at 1 Hz of Quartz Acoustic Resonators as a Function of Unloaded Q.

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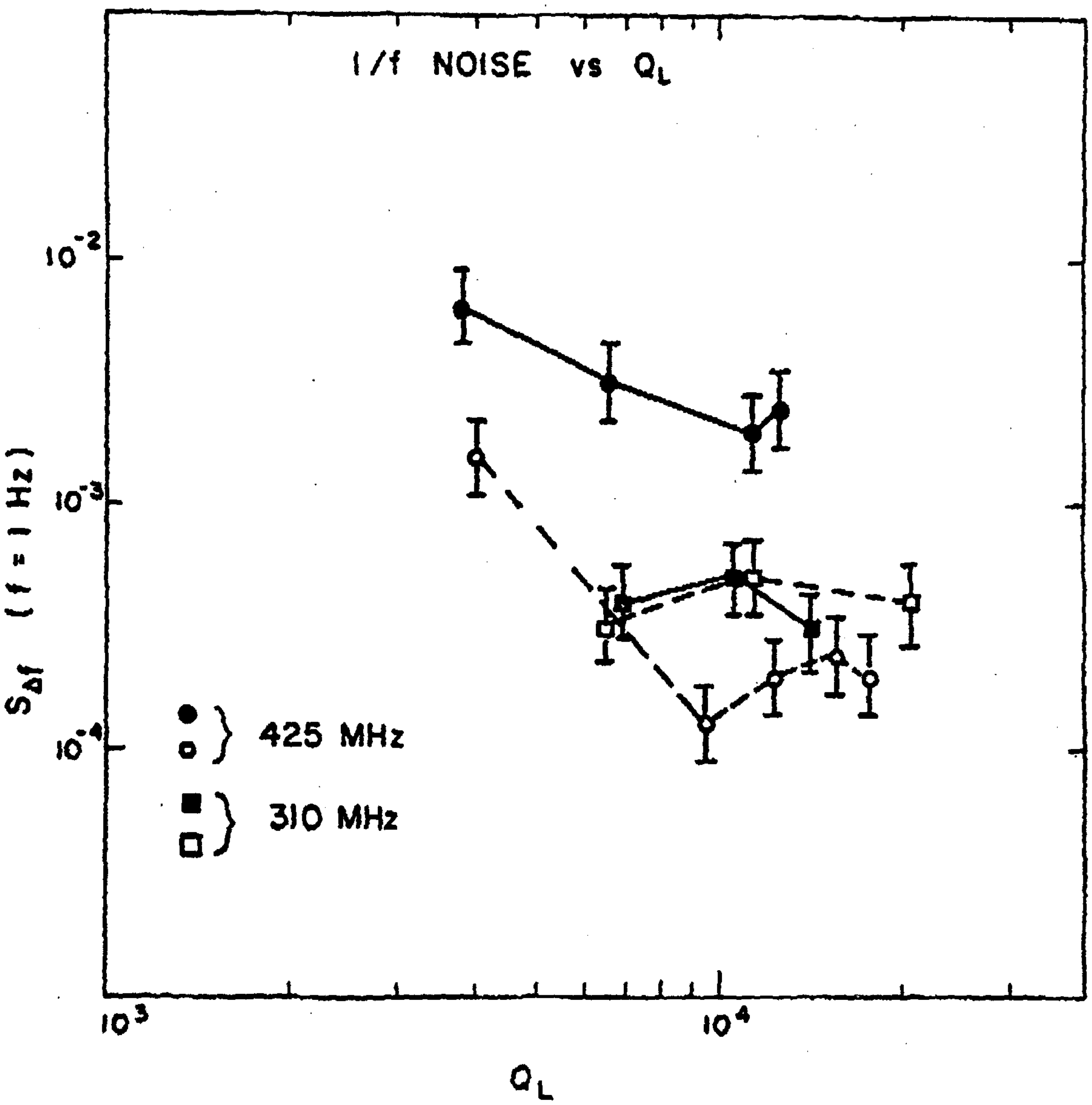
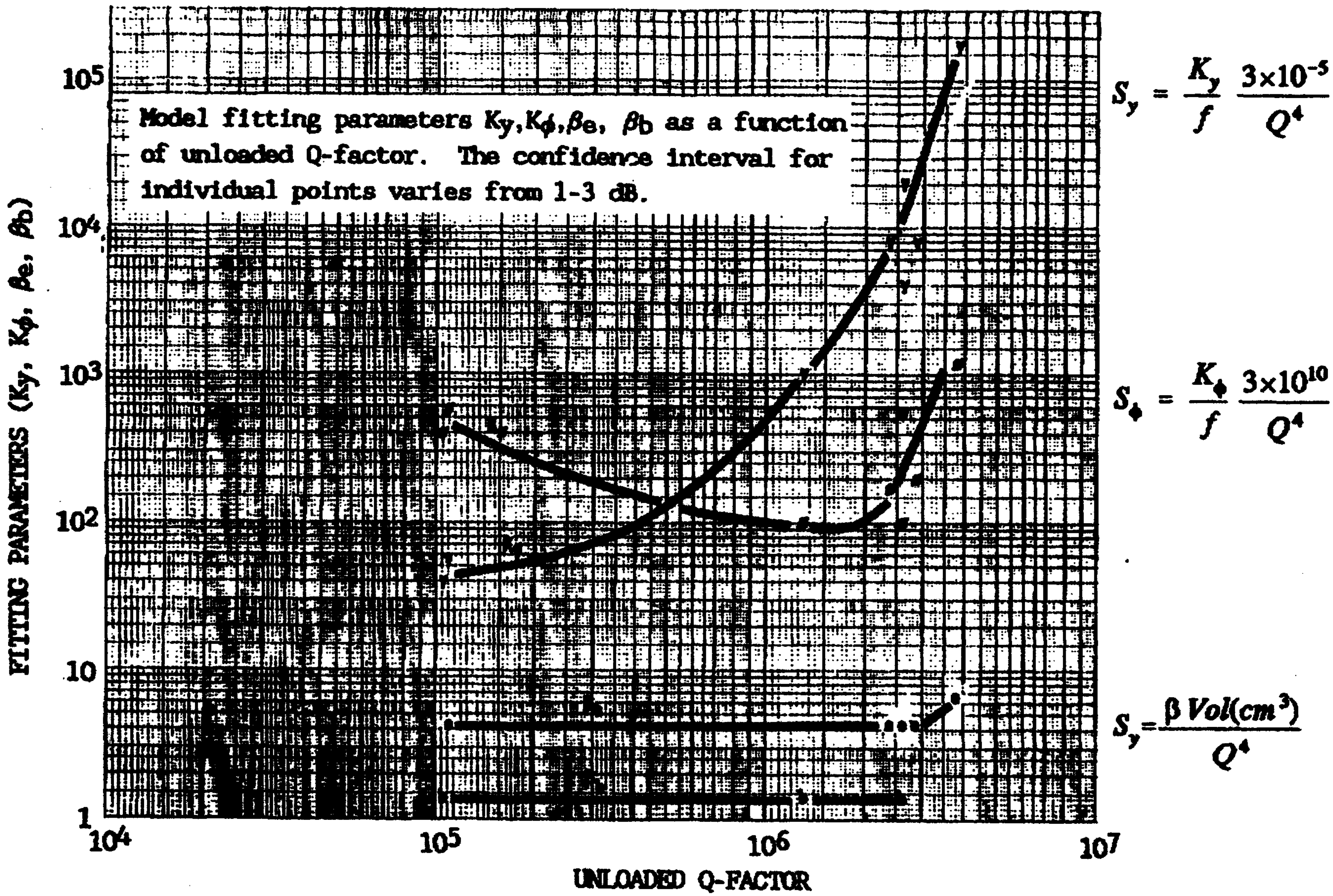


Figure 3 1/f Noise Level at 1 Hz of Four SAW Resonators as a Function of Loaded Q.



Model fitting parameters $K_y, K_\phi, \beta_e, \beta_b$ as a function of unloaded Q-factor. The confidence interval for individual points varies from 1-3 dB.
FITTING PARAMETERS (K_y, K_ϕ, β_e, β_b)
1
10
10^2
10^3
10^4
10^5
10^4
10^5
10^6
10^7
UNLOADED Q-FACTOR
$S_y = \frac{K_y}{f} \frac{3\times10^{-5}}{Q^4}$
$S_\phi = \frac{K_\phi}{f} \frac{3\times10^{10}}{Q^4}$
$S_y = \frac{\beta\, Vol(cm^3)}{Q^4}$

Fig 4

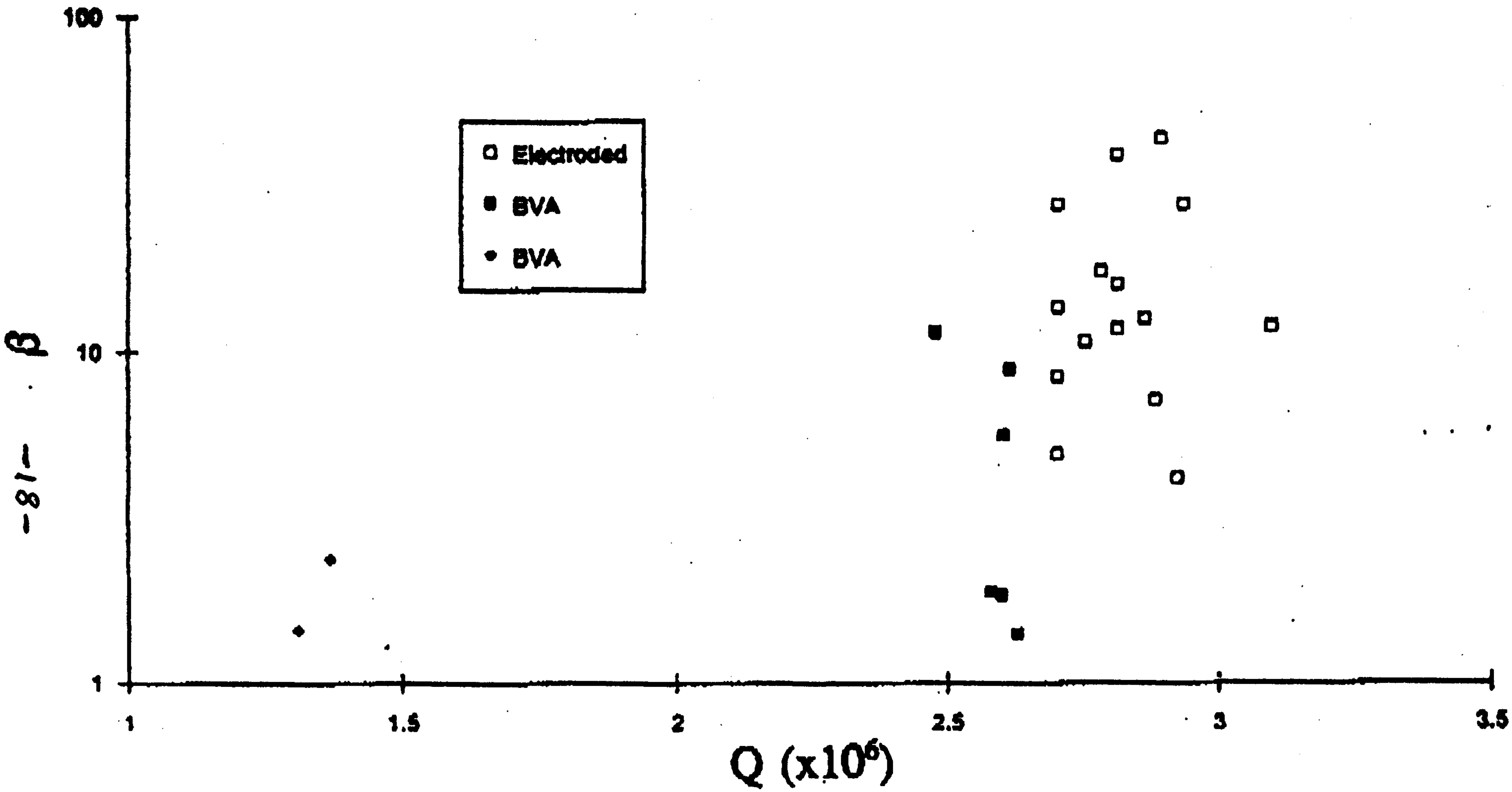


Figure 5. Fitting parameter β as a function of unloaded Q-factor for three types of resonators. The resonators in each group were matched in all known electrical parameters except Q-factor and 1/f noise [51].

be neglected and the pair can operate at virtually any frequency [45]. The oscillator approach makes it possible to compare many different resonators one at a time. The noise of individual oscillators can be derived by measuring the phase noise between 3 oscillators [48].

Figure 3 taken from [38] is one of several studies showing that the 1/f level is virtually independent of the loaded Q-factor. This is in complete agreement with the theoretical model. In practical oscillators there is a dependence on loaded Q-factor only when the phase noise of the sustaining electronics contributes to the overall noise level.

We have analyzed 1/f frequency noise as a function of unloaded Q, volume under the electrodes, and frequency. For a given resonator geometry and manufacturer, we have taken the best values of $S_y(f)$ reported in an attempt to remove the effects of poor crystals or electronics. In Fig. 4 we have taken all of the precision data available with unloaded Q-factor, electrode volume, and frequency stability and plotted them according to the three models. Except for the 2.5 MHz resonator where $Q\nu_0 = 0.95\times10^{13}$, the $Q\nu_0$ product for all resonators is near 1.2×10^{13} (this is close to the material limit for AT and SC cut resonators). The curve labeled K_y shows the fit of the data to the model [37] $S_y(f) = K_y/f(3 \times 10^{-5}/Q^4)$. K_y varies about a factor of 500 for Q-factors between 10^5 and 3.8×10^6 (resonator frequencies between 2.5 and 100 MHz).

The curve labeled K_ϕ shows the fit of the same resonator data to the model [38] $S(f) = K_\phi/f^3 (3 \times 10^{10}/Q^4)$. K_ϕ varies about a factor of 10 for the same range in Q-factor. Curves β_e and β_b show the fit of the same resonator data to the model $S_y(f) = \beta/f\ (\mathrm{Vol}/Q^4)$, where β_e is for SC and AT resonators with electrodes plated on the resonator and β_b is for BVA-style AT and SC resonators [43]. Volume between the electrodes (in cm^3) is used to approximate the volume of quartz contributing to the output power. <u>The β factors are remarkably constant for Q-factors from 10^5 to 3.8×10^6</u>.

Figure 5 shows the dependence of β_e on Q-factor for a number of electroded resonators of the same type from a single manufacturer for 3 resonator types as measured by Norton [44,51]. The wide variation in β_e for the same style resonator and Q-factor indicates that acoustic loss is not the only mechanism contributing to the noise level. The data for this graph was taken from measurements of (100 s) and may have been biased high by random walk FM noise in some resonators.

IV. 4. Discussion

The analysis of the most stable quartz resonators indicates that the 1/f frequency noise level depends on volume between the electrodes and unloaded Q-factor in relatively good agreement with eq. (30) considering the fact that the estimation of $<\omega>$ and N is only approximative. The nonadjustable parameters β_e and β_b are virtually constant versus unloaded Q-factor, which is in stark contrast to fitting parameters K_ϕ and K_y. It is not surprising to us that β_e and β_b are different for the two types of resonators, since energy trapping and electrode stress are considerably different. Fig. 5 shows that there are other noise processes besides acoustic losses that affect the 1/f noise level in some resonators.

Although we have analyzed only the data for a few resonators, the consistency of β_e and β_b over a factor of 40 in Q-factor and resonator frequency and the general agreement for the magnitude of β between theory and experiment give us some confidence that this new model can be used to predict the best performance of different resonator geometries.

This new volume model predicts that a resonator having smaller electrodes would have a lower level of 1/f frequency noise than another one with the same frequency and Q-factor but larger diameter electrodes. The decrease in electrode area would increase the impedance levels and degrade the wideband noise somewhat. For most resonators the wideband noise is dominated by the electronics and not the resonator. The increase in series resistance, obtained by decreasing the electrode area by a factor of 4, would probably be tolerable from the standpoint of wideband noise but might require a change in loop gain.

BT resonators are potentially interesting in that they offer a $Q\nu_0$ product approximately three times higher than that of AT and SC resonators. BT cuts are roughly as sensitive to temperature transients as AT cuts. Therefore to achieve parts in 10^{-14} frequency stability with BT cuts would require temperature stabilities of order 10^{-9} K/s or 100 times better than is required for SC cut resonators [51].

Based on these early observations it appears that the level of 1/f frequency noise in quartz may yet be improved to the low 10^{-14} level by applying one or more of the following techniques: reducing the electrode area, using BVA type resonators, going to lower frequencies, using BT cut resonators. It must be remembered that acceleration induced effects become more dominant as the stability improves.

V. PRACTICAL APPLICATION TO QUANTUM 1/f NOISE IN INFRARED DETECTORS

Quantum 1/f Noise is a fundamental aspect of quantum mechanics, representing universal fluctuations of physical process rates R and cross sections σ given by the fractional (or relative) spectral density $S(f) = 2\alpha A/fN$. Therefore it is present in the process rates generating the dark current observed in junction photodetectors, such as *diffusion* (scattering cross sections fluctuate) in diffusion-limited junctions, and *recombination* in the recombination-limited regime. One is therefore tempted to expect similar fluctuations in the *photogeneration* of electron-hole pairs. However, as we show below, the corresponding quantum 1/f coefficient is zero, precluding the existence of quantum 1/f fluctuations in the photogeneration rate. Here N is the number of carriers used to define or measure the process rate or cross section considered.

For an arbitrary process involving a total of n incoming and outgoing charged particles, the nonrelativistic quantum 1/f coefficient is given [52] by

$$2\alpha A = (4\alpha/3\pi c^2) \sum_{i;j=1}^{n} \eta_i \eta_j q_i q_j (v_i - v_j)^2, \tag{32}$$

where the summation runs over the charges q_i and velocities v_i of all incoming ($\eta_i=-1$) and outgoing ($\eta_i=1$) particles (altogether n of them) in the process whose quantum 1/f noise we want to find, and α is Sommerfeld's fine-structure constant, $e^2/\hbar c=1/137$. In a photoelectric process a photon ($q=0$) is absorbed, and a pair of oppositely charged particles is generated ($\eta=1$) with velocities v_1 and v_2 which are either zero, or quickly decay to zero in a time negligible with respect to the reciprocal frequency at which we calculate the quantum 1/f noise. Thus in our case there are no incoming charged particles, and $n=0+2=2$. The αA coefficient of a photogeneration process is therefore zero,

$$\alpha A_{ph} = (1,1)+(2,2)+(1,2)+(2,1) = 0+0+ (4\alpha/3\pi c^2)(v_1-v_2)^2 = 0. \tag{33}$$

All photogenerated carriers of the right sign are collected in the well of the charge-coupled device, although they may generate quantum 1/f voltage fluctuations on their way. Since usually only the number of carriers collected at read-out matters, no quantum 1/f noise will be observed in a photoelectric CCD as long as the dark current is negligible with respect to the photocurrent. This is in agreement with the experiments performed by Mooney [53] of RADC-

Hanscom AFB. The same considerations apply to Metal-Insulator-Semiconductor (MIS) photodetectors [54].

VI. QUANTUM 1/f MOBILITY FLUCTUATIONS IN HOMOGENEOUS SEMICONDUCTOR SAMPLES

VI. 1. Introduction

A first principles calculation of quantum 1/f cross correlations performed [55] for the first time in 1987 by one of us [56] has yielded a slightly different result compared to earlier expectations. This same new form of the quantum 1/f cross correlations was derived also with a different method by Van Vliet in 1989. It differs from the old form used in the 1985 calculation of Kousik et. al. by a correction which is zero when the momentum changes of the two current carriers involved in the cross correlation are identical, but increases to finite values when the momentum differences caused by the scattering process are different. The correction is proportional to the squared difference of the two momentum changes. We have repeated all calculations in the original paper by Kousik et.al. [56], obtaining both for impurity scattering and for the various types of phonon scattering new analytical expressions which show a considerable increase of the final quantum 1/f noise. The results obtained are applicable both to direct and indirect bandgap semiconductors.

We have performed an analytical calculation of mobility fluctuations in silicon and gallium arsenide, using the new quantum 1/f cross-correlations formula. This calculation is of major importance for the 1/f noise-related optimization of the two types of materials, and of the many devices constructed with them for military and civilian applications in the electronic and opto-electronic industry.

The new cross-correlation formula gives the cross-spectral density which describes the way in which simultaneous quantum 1/f scattering rate fluctuations ΔW observed in the direction of the outgoing scattered wave-vector K' are correlated with those in the K'' direction, when the two corresponding incoming current carriers have the wave vectors K_1 and K_2:

$$S_{\Delta W}(K_1,K';K_2,K'';f) = (2\alpha/3\pi f)(\hbar/m^* c)^2 W_{K_1,K'} W_{K_2,K''}[(K'-K_1)^2+(K''-K_2)^2]\delta_{K_1,K_2}. \qquad (34)$$

The form conjectured by us earlier had $2(K'-K_1)(K''-K_2)$ in place of the rectangular bracket. The difference between the rectangular bracket and $2(K'-K_1)(K''-K_2)$ is the perfect square $[(K'-K_1)-(K''-K_2)]^2$. Therefore we expect the new results to be always larger than the results obtained on the basis of the previously conjectured form.

VI. 2. Impurity Scattering

For impurity scattering of electrons in solids, fluctuations $\Delta\tau$ of the collision times τ will cause mobility fluctuations

$$\Delta\mu_{band}(t) = [e/m^*<<v^2>>]\Sigma_K v_K{}^2 \Delta\tau(t) n_K, \tag{35}$$

where $<<v^2>>$ is both the average over all states of wave-vectors K, with occupation numbers n_K, in the conduction band, and the thermal equilibrium average of the quadratic carrier velocities. With the help of the relation

$$1/\tau(K) = (V/8\pi^3)\int(1 - \cos\theta'/\cos\theta) W_{K,K'} d^3K', \tag{36}$$

the mobility fluctuations are reduced to fluctuations of the elementary scattering rates $W_{K,K'}$, governed by Eq. (34). Here V is the volume of the normalization box which disappears in the final result, θ and θ' respectively the angles K and K' form with the direction of the applied field. One finally obtains after tedious multiple integrations

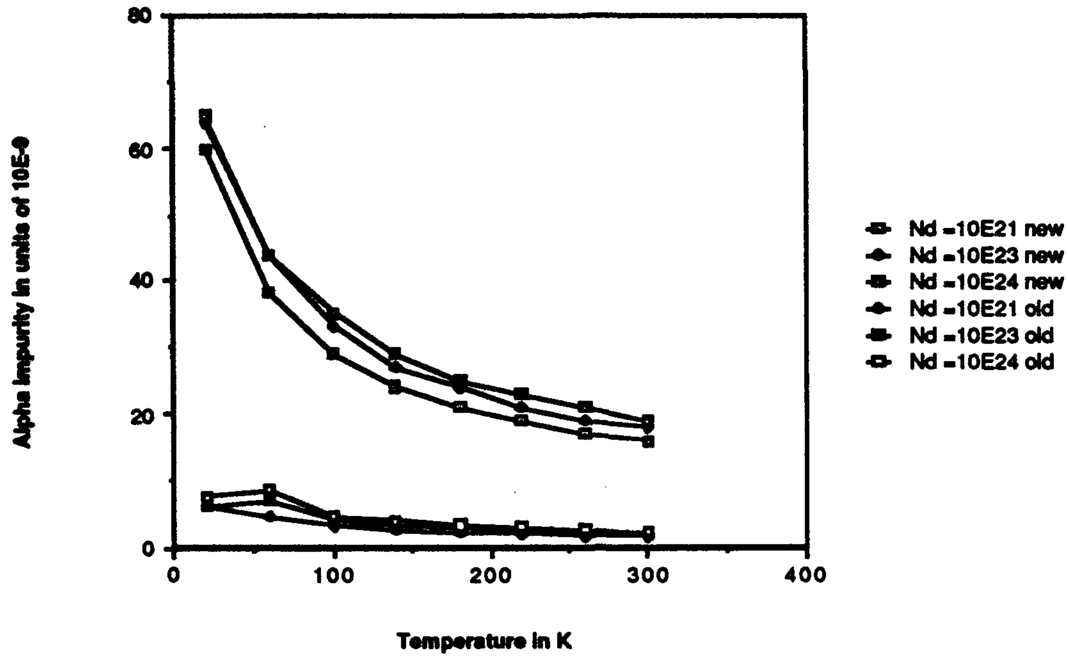
$$\mu^{-2}S_{\Delta\mu}(f) = [256\pi\alpha\kappa^2\varepsilon^4\hbar^{12}/3m^{*8}Z^4e^8N_i{}^2](1/f)\Sigma_K K^{10}[\ln(1+a^2)-a^2/(1+a^2)]^{-3}[(2a^2+a^4)/(1+a^2)-2\ln(1+a^2)]F(E_K)[\Sigma_K v_K{}^2\tau(K)F(E_K)]^{-2}, \tag{37}$$

where $a=2K/\kappa$, $\kappa^2=e^2n(T)/\varepsilon k_BT$, $n(T)$ is the electron concentration, $F(E_K)=\exp(E_F-E_K)$ for non-degenerate semiconductors, N_i the concentration of impurities of charge Ze and ε the dielectric constant. The corresponding partial Hooge parameter for impurity scattering is thus

$$\alpha_i = [4\sqrt{2}\pi\alpha\kappa\hbar^5N_c/3m^{*7/2}(k_BT)^{3/2}c^2]\int_0^\infty dx\, x^{11/2}e^{-x}$$
$$[\ln(bx+1)-bx/(bx+1)]^{-3}[(2bx+b^2x^2)/(bx+1)-2\ln(bx+1)]$$
$$\{\int_0^\infty dx\, x^3e^{-x}[\ln(bx+1)-bx/(bx+1)]^{-1}\}^{-2}. \tag{38}$$

This result is graphed below for three different values of the donor concentration N_d and is compared with old results obtained by simply recalculating the old analytical expression [56]

Alpha Impurity New and Old



As expected, the new cross correlation formula leads to higher α_i values than the previously conjectured expression. This was mentioned in connection with Eq. (34) above.

VI. 3. Acoustic Electron-Phonon Scattering

In this case the calculation is similar, and leads to the result

$$\alpha_{ac} = [32\pi\alpha N_c m^* C^7 \hbar^3/3c^2 k_B T)^4]\{(1/R^2)\int_1^\infty dx\, x^{-4}$$

$$[(x-1)^7/7+(R+1)(x-1)^6/6+R(x-1)^5/5]$$

$$[(x-1)^5/5+(R+1)(x-1)^4/4+R(x-1)^3/3]\exp(-x^2/4R)$$

$$+\int_0^1 dx\, x^{-4}[(x+1)^5/5-(x+1)^6/6+(x-1)^5/5+(x-1)^6/6]$$

$$[(x+1)^3/3+(x-1)^4/4+(x-1)^3/3-(x+1)^4/4]\exp(-x^2/4R)$$

$$+\int_1^\infty dx\, x^{-4}[(x+1)^5/5-(x+1)^6/6][(x+1)^3/3-(x+1)^4/4]\exp(-x^2/4R)\}, \qquad (39)$$

where $R=k_BT/2m^*C_1^2$, C_1 is the deformation potential, and N_c is the effective density of states for the conduction band.

VI. 4. Non-Polar Optical Phonon Scattering

This time one obtains

$$\alpha_{n.o.ph} = [8\pi\sqrt{2\hbar\omega_o\alpha}N_c\hbar^2/3m^{*5/2}c^2\omega_o]\{\int_0^\infty dx x^{5/2}$$
$$[(F+1)(x-1)^{1/2}\theta(x-1)+F(x+1)^{1/2}]^{-4}$$
$$[(F+1)^2(x-1)(2x-1)\theta(x-1)+F^2(x+1)(2x+1)]\exp(-\hbar\omega_o x/k_BT)\}$$
$$\{\int_0^\infty dx x^{3/2}[(F+1)(x-1)^{1/2}\theta(x-1)+F(x+1)^{1/2}]^{-1}\exp(-\hbar\omega_o x/k_BT)\}^{-2}, \qquad (40)$$

where $F=[\exp(\hbar\omega_o/k_BT)-1]^{-1}$, and ω_o is the optical phonon frequency.

VI. 5. Polar Optical Phonon Scattering

Proceeding as in Secs. 2 and 4, we obtain

$$\alpha_{p.o.ph} = [8\pi\sqrt{2\hbar\omega_l\alpha}N_c\hbar^2/3m^{*5/2}c^2\omega_l]\{\int_0^\infty dx x^4$$
$$[F^2(x+1)^{1/2}\ln(2x^{1/2}+2(x+1)^{1/2})$$
$$+(F+1)^2(x-1)^{1/2}\ln(2x^{1/2}+(x-1)^{1/2})\theta(x-1)]\exp(-\hbar\omega_l x/k_BT)$$
$$[(F+1)\mathrm{arcsinh}(x-1)^{1/2}\theta(x-1)+F\mathrm{arcsinh}(x^{1/2})]^{-4}\}. \qquad (41)$$

Here ω_l is the longitudinal phonon frequency.

VI. 6. Intervalley Scattering

This type of scattering, present in indirect bandgap semiconductors, transfers electrons from one of the six minima (or valleys) of the conduction band energy in k-space to one of the other five minima. Transitions between a valley and the nearest valley, which is along the same k-space direction in the next copy of the first Brilloin zone in the periodic zone scheme, are of the Umklapp type, and are called g-processes. Transitions to the four valleys present in the same zone along the other two k-space directions are called f-processes. Repeating a previous calculation [56] on the basis of the new cross-correlation formula (34), we obtain fo g-processes

$$\alpha_g = [8\pi\sqrt{2\hbar\omega_{ij}\alpha}N_c\hbar^2/3m^{*5/2}c^2\omega_{ij}]\{\int_0^\infty dx x^{5/2}$$
$$[(F+1)(x-1)^{1/2}\theta(x-1)+F(x+1)^{1/2}]^{-4}$$
$$[(F+1)^2(x-1)(2x-1)\theta(x-1)+F^2(x+1)(2x+1)]\exp(-\hbar\omega_{ij}x/k_BT)\}$$

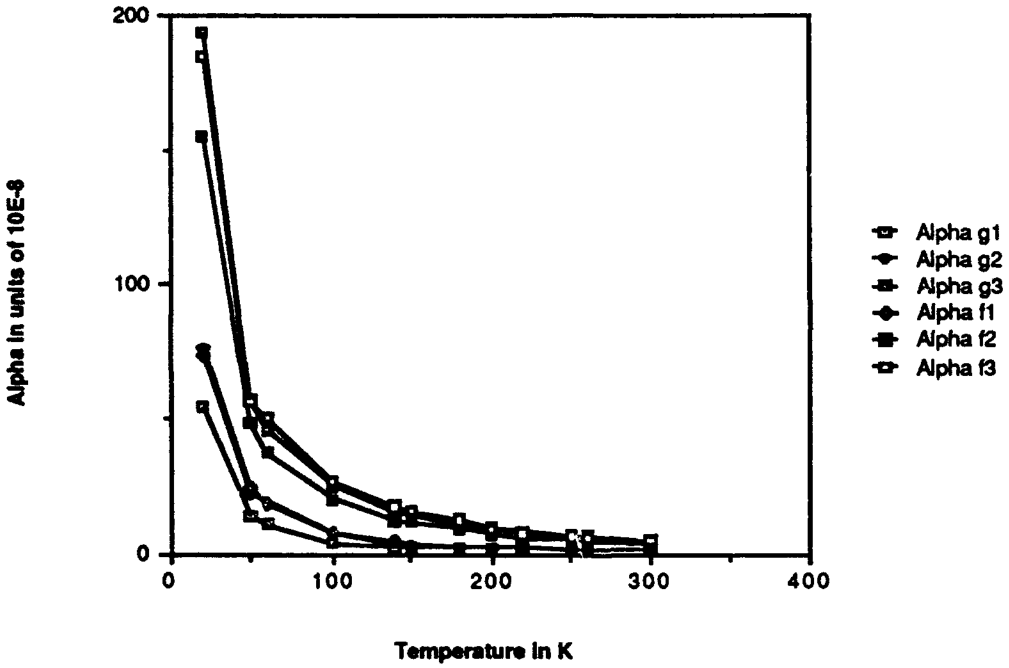
$$\left\{ \int_0^\infty dx\, x^{3/2}[(F+1)(x-1)^{1/2}\theta(x-1)+F(x+1)^{1/2}]^{-1} \exp(-\hbar\omega_{ij}x/k_BT)\right\}^{-2}, \quad (42)$$

where $\hbar\omega_{ij}$ is the phonon energy corresponding to the momentum difference required by the intervalley transition. For the corresponding f-process we obtain [57]

$$\alpha_f = (k_o/q_o)^2\alpha_g, \quad (43)$$

where k_o/q_o is the ratio between the position vector of a conduction band energy minimum in k space, and twice the distance of the minimum from the Brillouin zone boundary, 0.85/0.3 for silicon. There are three g-type alphas α_{g1}, α_{g2} and α_{g3} (from LA, TA and LO phonons repectively) and three f-type contributions α_{f1}, α_{f2} and α_{f3} (from TA, LA and TO phonons). Their values are given in the graph below and are a few times larger than the old values.

Alpha for Intervalley Scattering Processes



The various quantum 1/f contributions derived here can be approximately superposed to yield the resultant quantum 1/f coefficient according to the rule

$$\alpha_H = \Sigma_i(\mu/\mu_i)^2\, \alpha_i \quad (44)$$

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VIII.PAPERS PUBLISHED DURING THIS GRANT PERIOD

A. Published in refereed journals and conference proceedings which were subject to peer review.

P.H. Handel: "1/f Noise Criterion for Chaos in Nonlinear Systems", Proc. XI. International Conference on Noise in Physical Systems and 1/f Fluctuations, Nov. 24-27, 1991, Kyoto, Japan (Invited Paper), T. Musha, S. Sato and M. Yamamoto Editors, Ohmsha Ltd. Publishing Co., 3-1 Kanda Nishi-cho, Chiyoda-ku, Tokyo 101, Japan, pp.151-157.

P.H. Handel and E. Bernardi: "Ab Initio Derivation of Coherent Quantum 1/f Chaos", Ibid. pp.159-162

T.H. Chung, P.H. Handel and J. Xu: "Quantum 1/f Mobility Fluctuations in Semiconductors", Ibid. pp. 163-166.

A'. Published in conference proceedings which were not subject to peer review: five talks at the APS March 1992 Meeting in Indianapolis.

P. Handel, "Fundamental 1/f Noise and Quantum 1/f Noise in Quartz", APS Bull. 37, 267 (1992)

T. Chung, P. Handel and J. Xu, "Quantum 1/f Mobility and Diffusion Fluctuations", p. 267 ibid.

J. Xu and P.H. Handel, "Second Quantized Derivation of Quantum 1/f Noise in the presence of a Thermal Background", ibid., p. 268.

Xuewei Hu and P. H. Handel, "Fractal Dimension of Quantum 1/f Noise", ibid. p. 268.

E. Bernardi and P.H. Handel, "Derivation of Coherent Quantum 1/f Chaos", ibid., p. 268

P. H. Handel, "Sufficient Criterion for a 1/f Spectrum of Chaos in Nonlinear Systems", ibid., p. 269.

B. In Press (accepted for publication)

F.L. Walls, P.H. Handel, R. Besson and J.J. Gagnepain: "A New Model of 1/f Noise in BAW Quartz Resonators", Proc. of the 46th Ann. Frequency Control Symp., 1992.

C. Papers submitted for publication, but not yet accepted

P.H. Handel: "Quantum 1/f Fluctuation of Physical Cross Sections and Process Rates", p.1-9, Submitted to J of Physics C.

P.H. Handel: "Sufficient Criterion for a 1/f Spectrum of Chaos in Nonlinear Systems" submitted to the Physical Review.

P.H. Handel: "Coherent Quantum 1/f Chaos Effect", submitted to Phys.Rev.Lett.

D. Invited Talks

"1/f Noise Criterion for Chaos in Nonlinear Systems" in the "Quantum 1/f Noise" session at the "XI Int. Conf. on Noise in Physical Systems and 1/f Noise" in Kyoto (Japan) Nov. 1991.

"Quantum 1/f Noise in Quartz Resonators" at NIST (Natl. Inst. of Standards), Boulder, Colorado, January 1992.

E. Service

Chaired the session on "1/f Noise and Quantum Chaos" at the APS (1992) "March Meeting" in Indianapolis.

Chaired a session at the "XI Int. Conf. on Noise in Physical Systems and 1/f Noise" in Kyoto (Japan) Nov. 1991.

Organizer and Chair of the 5th van der Ziel Symposium on Quantum 1/f Noise and Other Low-Frequency Fluctuations", and of a "Quantum 1/f Retreat" at UMSL on May 22-25, 1992.

Chair and Organizer of the "XII Int. Conf. on Noise in Physical Systems and 1/f Noise" in St. Louis, MO, Aug 16-20, 1993.

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